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Flat Potential for Inflaton with a Discrete R-invariance in Supergravity

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Abstract

We show that a very flat potential of inflaton required for a sufficient inflation is naturally obtained in $N = 1$ supergravity by imposing a discrete R-invariance Z_n. Several cosmological constraints on parameters in the inflaton superpotential are derived. The reheating temperature turns out to be $(1 - 10^8)$GeV for the cases of n=3–10. Baryogenesis in this model is also discussed briefly.

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1. Introduction

Much effort has been done in building realistic particle-physics models based on $N = 1$ supergravity [1]. Although these theories provide a natural framework for producing soft supersymmetry (SUSY) breaking terms at low energies, there are potential cosmological problems. One of them is a difficulty to construct inflationary scenarios of the universe. Namely, to generate a sufficiently large expansion of the universe one must require an extreme fine tuning of several parameters making a very flat potential for the inflaton [2]. However, there has not been found any natural explanation on the existence of such a flat potential.

In this paper, we show that a discrete R-symmetry Z_n automatically leads to a flat potential at the origin of the inflaton field ϕ as long as the minimum Kähler potential $K(\phi, \phi^*) = \phi\phi^*$ is used. We derive cosmological constraints on parameters in the inflaton superpotential.

In this model, the inflaton superpotential does not vanish at the minimum of the inflaton scalar potential and hence the inflaton ϕ plays a role of the Polonyi field for generating a gravitino mass [3]. Together with the cosmological constraints, all relevant parameters in the superpotential are fixed such that the gravitino mass $m_{3/2}$ lies in 100GeV – 10TeV. With the obtained parameters in the superpotential we argue that the most natural choice for the discrete R-symmetry is Z_4, in which the reheating temperature is $T_R \sim O(1-100)$TeV depending on the gravitino mass $m_{3/2}$=100GeV – 10TeV. We give a brief comment on a possible scenario for baryogenesis in such a low temperature universe. Particle-physics problems related to the SUSY breaking in this model are also discussed.

2. Discrete R-symmetries and the inflaton potential

Let us define a discrete Z_n R-transformation on the inflaton field $\phi(x,\theta)$ as

$$\phi(x,\theta) \to e^{-i\alpha}\phi(x, e^{i\alpha/2}\theta), \quad \alpha = \frac{2\pi k}{n}, \tag{1}$$

with $k = 0, \pm 1, \pm 2 \cdots$. Assuming positive power expansion of a Kähler potential $K(\phi, \phi^*)$ and a superpotential $W(\phi)$ for the inflaton field ϕ, the general forms which

have the Z_n R-invariance are given by

$$K(\phi,\phi^*) = \sum_{m=1}^{\infty} a_m (\phi\phi^*)^m, \tag{2}$$

and

$$W(\phi) = \phi \sum_{l=0}^{\infty} b_l \phi^{ln}. \tag{3}$$

We take the minimum Kähler potential (a_1=1, other a_i=0) for simplicity. (We have found that the scenario described below does work with more general Kähler potential in eq.(2), as far as the coefficients $|a_2| < 3/8n(n+1)$ and $|a_i| \sim O(1)(i \geq 3)$.)

In the $N = 1$ supergravity [1], a scalar potential $V(\phi)$ is written as

$$V(\phi) = \exp\left(\frac{K(\phi,\phi^*)}{M^2}\right)\left\{(K^{-1})_{\phi}^{\phi^*}(D_\phi W)(D_\phi W)^* - \frac{3|W|^2}{M^2}\right\}, \tag{4}$$

with

$$\begin{aligned} D_\phi W &\equiv \frac{\partial W}{\partial \phi} + \frac{\partial K}{\partial \phi}\frac{W}{M^2}, & (5)\\ (K^{-1})_{\phi}^{\phi^*} &\equiv \left(\frac{\partial^2 K}{\partial\phi\partial\phi^*}\right)^{-1}, & (6)\end{aligned}$$

where M is the gravitational mass $M = M_{Planck}/\sqrt{8\pi} \simeq 2.4 \times 10^{18}$GeV. With the general form of $W(\phi)$ in eq.(3), we find the inflaton potential $V(\phi)$ has automatically flat directions at the origin. That is, the condition

$$\frac{\partial V}{\partial \phi} = \frac{\partial^2 V}{\partial\phi\partial\phi^*} = 0, \tag{7}$$

is always satisfied at the origin $\phi = 0$. Notice that the mass term $\phi\phi^*$ from the $\exp(K/M^2)$ is cancelled out by that in the curly bracket in eq.(4). This cancellation is very important for the inflaton ϕ to do its required job, and in all previous works [2] this cancellation is achieved only by a fine tuning of parameters in the superpotential $W(\phi)$.[1]

[1]There has been proposed an alternative approach, in which non-minimum Kähler potentials are used [4]. In this class of models, a fine tuning is necessary in the Kähler potential to achieve a very flat potential for the inflation.

To find the approximate form of $V(\phi)$ near the origin $\phi \sim 0$ we write the superpotential as

$$W(\phi) = \left(\frac{\lambda}{v^{n-2}}\right)\left(v^n\phi - \frac{1}{n+1}\phi^{n+1}\right) + \cdots. \tag{8}$$

Here, $\cdots$ represents higher power parts (ϕ^{kn+1} with $k \geq 2$). Since these parts are all irrelevant to the present analysis, we neglect them, hereafter. We easily see that the inflaton potential (4) near $\phi \sim 0$ is well approximated by

$$V(\phi) \simeq \left(\frac{\lambda}{v^{n-2}}\right)^2 \left\{ v^{2n} + \frac{1}{2}v^{2n}\left(\frac{|\phi|^2}{M^2}\right)^2 - v^n(\phi^n + \phi^{*n}) \right\}. \tag{9}$$

Clearly, n=2 does not give a flat potential and hence we do not consider the n=2 case. For n=3 and 4, we can neglect the $(|\phi|^2/M^2)^2$ term, since $v \ll M$ as we will see later. For $n \geq 5$, we also neglect this term assuming that the initial amplitude of ϕ is greater than $v(v/M)^{4/(n-4)}$.[2] Thus, in any case of $n \geq 3$ we neglect the $(|\phi|^2/M^2)^2$ term in analyzing the inflaton dynamics. In Fig. 1 we depict the exact $V(\phi)$ along the real axis Reϕ, which shows that the inflaton potential has always a very flat region near $\phi \sim 0$. We have also checked that $\phi \simeq v$ is the global minimum of the inflaton potential.

Whether the inflation occurs or not depends on initial conditions for the inflaton field ϕ. In the present model the initial amplitude of ϕ must be localized very near the origin in order to have a sufficiently large expansion of the universe. We simply assume ϕ and $\dot{\phi}$ satisfy the desired initial conditions $|\phi| \sim 0$ and $\dot{\phi} \sim 0$. We have no answer to a question of what physics the initial conditions for ϕ was set by, but if ϕ and $\dot{\phi}$ satisfy the initial condition, the maximum inflation takes place and this exponentially expanding part of the universe dominates the others.

Setting the phase of ϕ to vanish,[3] we identify the inflaton field with φ where φ is a real component of ϕ ($\varphi \equiv \sqrt{2}\mathrm{Re}\phi$). Thus, the relevant potential for φ is now given by

$$V(\varphi) \simeq \tilde{\lambda}^2\tilde{v}^4 \left\{ 1 - 2\left(\frac{\varphi}{\tilde{v}}\right)^n \right\}, \tag{10}$$

[2]This assumption on the initial amplitude $\varphi_i (\equiv \sqrt{2}\mathrm{Re}\phi_i)$ is consistent with $\varphi_i < \varphi_N$(with $N \sim 40$) given in eq.(17) as far as $v \ll M$ and $n \geq 5$. This initial condition $\varphi_i < \varphi_{N\sim 40}$ is derived to have a sufficient inflation as we will explain later.

[3]Notice that the phase $\chi(x)$ ($\phi = \sqrt{2}\varphi e^{i\chi}$) has a positive mass at $\varphi \neq 0$. Thus if one chooses the initial value of $\chi(x) \sim 0$, it stays there during the inflation.

with

$$\tilde{\lambda} \equiv \frac{1}{2}\lambda, \quad \tilde{v} \equiv \sqrt{2}v. \tag{11}$$

3. Cosmological constraints

The equation of motion for φ in the expanding universe is given by

$$\ddot{\varphi} + 3H\dot{\varphi} + \frac{dV}{d\varphi} = 0, \tag{12}$$

where H is the Hubble expansion rate. During the slow rolling regime of φ ($\ddot{\varphi} \ll 3H\dot{\varphi}$) the energy density of the universe is dominated by the inflaton potential $V(\varphi \sim 0)$, which gives a nearly constant expansion rate

$$H^2 \simeq \frac{V(\varphi \sim 0)}{3M^2} \simeq \frac{\tilde{\lambda}^2 \tilde{v}^4}{3M^2}. \tag{13}$$

In this inflationary epoch, the $\ddot{\varphi}$ term in eq.(12) can be neglected so that

$$\dot{\varphi} \simeq -\frac{V'(\varphi)}{3H} \simeq \frac{2n\tilde{\lambda}M}{\sqrt{3}\tilde{v}^{n-2}}\varphi^{n-1}. \tag{14}$$

The slow rolling regime ends at φ_f,

$$\varphi_f^{n-2} \simeq \frac{3}{2n(n-1)}\frac{\tilde{v}^n}{M^2}. \tag{15}$$

The cosmic scale factor grows exponentially $\sim e^N$ till the end of the inflation. With the above approximation, the e-folding factor N between the time t_f and t_N is given by

$$N = H(t_f - t_N) \simeq \frac{n-1}{3(n-2)}\left\{\left(\frac{\varphi_N}{\varphi_f}\right)^{2-n} - 1\right\}, \tag{16}$$

where φ_f (φ_N) represents the amplitude of the field variable φ at the time t_f (t_N). For a large N, φ_N is given by

$$\varphi_N \simeq \{2Nn(n-2)\}^{1/(2-n)}\left(\frac{\tilde{v}}{M}\right)^{2/(n-2)}\tilde{v}. \tag{17}$$

To solve the flatness and horizon problems, a sufficiently large expansion of the universe is required during the inflation [5]. In our model, the Hubble radius of the present

universe crossed outside of the horizon $N \sim 40$ e-folds before the end of the inflation.[4] This suggests that the initial amplitude of φ should be smaller than φ_N with $N \sim 40$.[5]

During the de Sitter phase, the density perturbation $(\delta\rho/\rho)$ arises from quantum fluctuations [6] of the inflaton field φ. It is roughly given by

$$\left(\frac{\delta\rho}{\rho}\right)_N \simeq \frac{3}{5\pi}\frac{H^3}{|V'(\varphi_N)|} \simeq \left(\frac{\tilde{\lambda}\tilde{v}^3}{10\sqrt{3}\pi n M^3}\right)\left\{\frac{\tilde{v}^2}{2Nn(n-2)M^2}\right\}^{(1-n)/(n-2)}. \tag{18}$$

The relation between the density perturbation and the quadrupole of the temperature fluctuation of cosmic microwave background (CMB) $\sqrt{\langle a_2^2\rangle}$ is given by [7]

$$\sqrt{\langle a_2^2\rangle} = \sqrt{\frac{5\pi}{12}}\left(\frac{\delta\rho}{\rho}\right)_{N\simeq 40}. \tag{19}$$

From the data on anisotropy of the CMB,

$$\left(\frac{\delta T}{T}\right) \simeq \sqrt{\frac{\langle a_2^2\rangle}{4\pi}} \simeq 6\times 10^{-6}, \tag{20}$$

observed by COBE [8], we derive a constraint

$$\left.\frac{\tilde{\lambda}\tilde{v}^3}{10\sqrt{3}\pi n M^3}\left\{\frac{\tilde{v}^2}{2Nn(n-2)M^2}\right\}^{(1-n)/(n-2)}\right|_{N\sim 40} \simeq 2\times 10^{-5}. \tag{21}$$

We are now at the point to discuss the SUSY breaking in the present model. Interesting is that the superpotential $W(\phi)$ does not vanish at the potential minimum $\phi \simeq v$,

$$W(\phi\simeq v) \simeq \frac{n}{\sqrt{2}(n+1)}\tilde{\lambda}\tilde{v}^3. \tag{22}$$

The gravitino mass $m_{3/2}$ is, then, given by

$$m_{3/2} \simeq e^{\langle K\rangle/2M^2}\frac{\langle W\rangle}{M^2} \simeq \frac{n}{\sqrt{2}(n+1)}\tilde{\lambda}\tilde{v}\left(\frac{\tilde{v}}{M}\right)^2. \tag{23}$$

[4]The reason why e-folding factor $N \sim 40$ is smaller than the usual value $N \sim 60$ is because the reheating temperature T_R in the present model is relatively low ($T_R \lesssim 100$TeV for $n = 3-5$). For $n \geq 6$, T_R is $(10^2 - 10^4)$TeV and in these cases, N becomes $N \sim 50$.

[5]This value of φ_N is much larger than the quantum fluctuation $\delta\varphi \sim H/2\pi$ unless λ is large $\lambda > (M/v)^{\frac{n-4}{n-2}}$. Furthermore, the change of φ in one expansion time is much larger than the quantum fluctuation. Thus, the evolution of φ can be discussed by solving the classical equation of motion in eq.(12).

Thus, the inflaton field ϕ is regarded as the Polonyi field for producing the gravitino mass [3].

From eqs.(21) and (23), we determine λ and v for a given sets of N and $m_{3/2}$. The results for $N = 40$ are shown in Table 1. For the obtained λ and v, we calculate the mass of ϕ as

$$m_\phi \simeq \lambda n v. \tag{24}$$

As seen in Table 1 the inflaton mass m_ϕ is predicted as,

$$m_\phi \simeq \left(10^7 - 10^{12}\right) \mathrm{GeV}. \tag{25}$$

This seems to contradict with the claim [9] that the mass of the Polonyi field is always at the gravitino mass scale $m_{3/2}$ ~100GeV – 10TeV. However, our result (25) is not inconsistent with their claim, since we have not demanded the cosmological constant to vanish. In fact, we have a non-zero cosmological constant at the inflaton potential minimum,

$$\Lambda_{\cos}^{\phi} = V(\phi \simeq v) \simeq -3m_{3/2}^2 M^2 \sim -(10^{10} - 10^{11}\mathrm{GeV})^4. \tag{26}$$

Therefore, we need to invoke some mechanism to cancel this negative cosmological constant. The simplest way is to introduce a U(1) gauge multiplet $V(\theta, x)$ in the hidden sector and add a Fayet-Iliopoulos D-term[6] [10, 11, 12], ξD, which shifts up the vacuum energy density by an amount of

$$\delta\Lambda_{\cos}^{D} = \frac{1}{2}\xi^2. \tag{27}$$

[6]In supergravity, the Fayet-Iliopoulos D-term can be written as (see Ref.[13] for notations)

$$\frac{3}{4}\int d^2\Theta M^2 \mathcal{E}\left(\bar{\mathcal{D}}\bar{\mathcal{D}} - 8R\right) \exp\left(-\frac{1}{3}\frac{\xi V}{M^2}\right).$$

Clearly, this term is not invariant under the U(1) gauge transformation $V \to V + \Lambda + \Lambda^\dagger$. Thus, we assume, here, that this D-term is induced by some mechanism for the U(1) breaking. Another solution to this problem may exist if the superpotential has a continuous R-symmetry [11, 12]. However, this symmetry conflicts with our case, since the superpotential (3) does not have the continuous R-symmetry. Therefore, we must also consider that the continuous R-symmetry is broken down to our discrete R-symmetry by some underlying physics at the Planck scale. A detailed argument on this problem will be given in future communication [14].

Thus we can always choose ξ so that the total cosmological constant vanish,

$$\Lambda_{\rm cos} = \Lambda^{\phi}_{\rm cos} + \delta\Lambda^{D}_{\rm cos} = 0. \tag{28}$$

This looks very artificial, but notice that a serious cosmological problem [15] associated with the light Polonyi field is not present due to the relatively large mass m_ϕ given in eq.(25). In any case, the presence of non-vanishing cosmological constant $\Lambda^{\phi}_{\rm cos}$ does not affect our inflation scenario, since $\Lambda^{\phi}_{\rm cos}$ in eq.(26) is always negligible[7] compared with the inflaton energy density $V(\phi \sim 0)$ in the inflationary epoch, as seen in Fig. 1.

Notice that the SUSY breaking due to the above D-term dominates over the F-term breaking by the inflaton ($F = D_\phi W$, see eq.(5)). Thus, the physical gravitino field ψ'_μ is composed mainly of the ψ_μ and λ through the super Higgs mechanism [1] as

$$\psi'_\mu = \psi_\mu + \frac{1}{6} e^{-\langle K\rangle/2M^2} M \frac{\langle D\rangle}{\langle W\rangle} \sigma_\mu \bar{\lambda}. \tag{29}$$

with

$$\langle D\rangle = \xi, \tag{30}$$

where ψ_μ and λ are the gravitino and the U(1) gaugino fields, respectively (in the two component Weyl representation). This physical gravitino field has a mass term

$$e^{\langle K\rangle/2M^2} \frac{\langle W\rangle}{M^2} \psi'_\mu \sigma^{\mu\nu} \psi'_\nu + h.c. \tag{31}$$

In the Minkowski space-time this gravitino mass becomes a physical pole mass of the gravitino propagator.[8] Furthermore, if we use the minimum Kähler potential for quark and lepton fields, the soft-SUSY breaking masses for squarks and sleptons come also from W in eq.(22).

Since the inflaton ϕ couples to the light observable sector very weakly (with strength $\sim (\lambda v^2/M^2)$), the decay width Γ_ϕ is very small as

$$\Gamma_\phi \sim \left(\frac{\lambda v^2}{M^2}\right)^2 m_\phi. \tag{32}$$

[7]Notice that we have not required, in our analysis, the cosmological constant to be negligibly small. The main reason why we have $\Lambda^{\phi}_{cos} \ll V(\phi \sim 0)$ is that the vacuum-expectation value v is very small, $(v/M) \sim 10^{-3} - 10^{-2}$.

[8]Notice that the presence of the gravitino mass given in eq.(23) does not mean the SUSY breaking in the anti-de Sitter space-time.

For the $n \leq 5$ case, this leads to a reheating temperature T_R [5],

$$T_R = O(10^{-2} - 10)\mathrm{GeV}, \tag{33}$$

which is too low to create enough baryon-number asymmetry of the universe. In the case of $n \geq 6$, the reheating temperature is $O(100)$GeV. But as we will see later, models with $n \geq 5$ have a physical cut-off smaller than the Planck scale. Therefore, if one requires the cut-off scale to be larger than the Planck scale, one should take n=3 or 4 and, hence a new interaction is necessary to get a sufficiently high reheating temperature for baryogenesis.

To have a faster decay of ϕ, we introduce a new singlet chiral supermultiplet $N(x,\theta)$ with a half Z_n charge of ϕ, that is

$$N(x,\theta) \to e^{-i\alpha/2} N(x, e^{i\alpha/2}\theta). \tag{34}$$

Then, we have a new interaction term in Kähler potential

$$K(\phi,\phi^*,N,N^*) = \frac{g}{M}\phi^* NN + h.c.\ , \tag{35}$$

and N can have a mass term

$$W(N) = \frac{m_N}{2}NN. \tag{36}$$

Provided $2m_N < m_\phi$, we have a much faster ϕ decay mode, $\phi \to NN$, with decay rate

$$\Gamma_{\phi\to NN} = \frac{g^2}{8\pi}\frac{m_\phi^3}{M^2}\left(1 - \frac{4m_N^2}{m_\phi^2}\right)^2. \tag{37}$$

With this decay width, we estimate the reheating temperature $T_R \sim (1 - 10^5)$TeV for n=4–10 (see Table 1).[9] Only for the case of n=3, the reheating temperature seems to be too low ($T_R \sim (1 - 100)$GeV) for the baryogenesis as discussed below.

It is a straightforward task to incorporate the Fukugita-Yanagida mechanism [18] for baryogenesis in the present model, identifying the singlet N with three families of

[9] In the recent article Fischler have derived a new constraint on the reheating temperature $T_R \lesssim (10^2 - 10^5)$GeV [16] to solve the gravitino problem [17]. If one adopts this constraint, there are left consistent only the n =3, 4 and 5 cases. However, it is not clear to us if this new constraint on T_R is relevant.

the right-handed neutrinos.[10] Taking Z_n charges for all quarks Q and leptons L to be the same as that of N and assuming Higgs multiplets H and $\bar{H}$ have zero-Z_n charges,[11] we find that the N's can have the standard Yukawa couplings

$$W_{\text{Yukawa}} = g_{ij} L_i N_j H, \tag{38}$$

where i and j represent the family indices. The decay $N \to LH$ produces a lepton-number asymmetry if g_{ij} has a CP violating phase. The produced lepton-number asymmetry can be converted to the baryon-number asymmetry [18] through the anomalous electroweak processes [21] at high temperature $T \gtrsim O(100\text{GeV})$. Therefore, the reheating temperature should be higher than $O(100\text{GeV})$.[12] As seen in Table 1 this condition is satisfied in models with $n \geq 4$, while in the $n = 3$ case the gravitino mass less than $O(10\text{TeV})$ is unlikely.[13]

4. Conclusion

Some comments are in order.

[10]Through the seesaw mechanism [19], the light neutrino mass m_ν is given by $m_\nu \simeq m_D^2/m_N$ with m_D being the Dirac mass of neutrino. If all $m_N \sim 10^9$GeV the MSW solution [20] to the solar ν problem suggests $m_D \simeq 0.1 - 0.01$GeV.

[11]The invariant mass term $W = \mu H\bar{H}$ is forbidden by the Z_n-symmetry. If one considers the exact Z_n-symmetry one needs a Yukawa interaction $W = h\phi H\bar{H}$ to produce the invariant mass. To give a week-scale mass for H and $\bar{H}$ one must choose h very small, $h \sim 10^{-14} - 10^{-12}$.

[12]If $2m_{N_1} < m_\phi < 2m_{N_{2,3}}$, only the N_1 is responsible for the baryogenesis. Suppose a hierarchy in the Yukawa coupling, $|g_{33}| \gg \text{other}|g_{ij}|$, one obtains the lepton-number asymmetry through the $N \to LH$ and L^*H^* decay processes as [18]

$$\Delta L/s \sim 10^{-5} \frac{1}{8\pi} |g_{33}|^2 \delta \left(\frac{m_{N_1}}{m_{N_3}} \right),$$

where 10^{-5} is a dilution factor due to the entropy production of N decays, and δ is the CP-violating phase of g_{ij}. This lepton-number asymmetry is converted to the baryon-number asymmetry and one gets [22]

$$\Delta B/s \simeq \frac{4n_f + 4}{12n_f + 13} \Delta L/s,$$

with n_f being the number of families ($n_f = 3$).

With this result, one may easily explain the observed baryon-number asymmetry $\Delta B/s \simeq 10^{-10} - 10^{-11}$, taking $g_{33} = O(1)$ and $\delta(m_{N_1}/m_{N_3}) \sim 10^{-4}$. Notice that the N_1 decay processes are always out of equilibrium since $m_{N_1} \gg T_R$ and hence we do not have any additional constraint (so-called out-of-equilibrium condition) on m_{N_1}.

[13]If one uses the Affleck-Dine mechanism [23] for the baryogenesis, the $n = 3$ case is not ruled out.

(i) Since the superpotential $W(\phi)$ contains a non-renormalizable term $\frac{\lambda}{v^{n-1}}\frac{1}{n+1}\phi^{n+1}$, the Born unitarity is violated in the process $\phi+\phi^* \to (n-1)\phi+(n-1)\phi^*$ above a very high-energy scale. Thus we need a physical cut-off scale Λ. By using a simple power counting, we estimate the cut-off scale to be

$$\Lambda \sim \lambda^{1/(1-n)} v. \tag{39}$$

If one imposes $\Lambda \geq M$ or M_{Planck}, one finds that only the n=3 and n=4 cases are consistent (see Table 1).

(ii) The gaugino masses come from the non-minimal kinetic term of gauge multiplet W_α [13],

$$\frac{1}{8g^2}\int d^2\Theta\mathcal{E} f(\phi) W^\alpha W_\alpha + h.c., \tag{40}$$

where f is the kinetic function for the gauge multiplet. Eq.(40) induces the gaugino mass term

$$\frac{1}{4}\left\langle \frac{\partial f}{\partial \phi}\ D_\phi W \right\rangle \tilde{g}^\alpha \tilde{g}_\alpha + h.c. \tag{41}$$

Suppose $f = \kappa\phi/M$, we get a very small gaugino mass[14]

$$m_{\tilde{g}} \simeq \frac{\kappa\lambda n v^4}{4(n+1)M^3} \sim \begin{cases} 10^{-3} m_{3/2}, & \text{for } n=4, \\ 10^{-2} m_{3/2}, & \text{for } n=3, \end{cases} \tag{42}$$

with $\kappa = O(1)$. Therefore, to obtain enough large gaugino masses[15] we must take account of radiative corrections from some heavy particles[16] [25, 26]. This situation is very much similar to in the scenario of dynamical SUSY breaking [27]. Whether the radiative corrections give rise to sufficiently large masses for the gauginos depends on detailed physics at the high-energy (perhaps the Planck) scale. Therefore, it may be safe to conclude that the gaugino masses are much smaller than those of the squarks and sleptons unless a huge number of heavy particles exists at the high-energy scale.

In conclusion, we have shown that a discrete R-symmetry Z_n naturally produces a very flat potential for the inflation. Particularly for the case of $n = 4$, the effective

[14] We thanks M.Yamaguchi for pointing out this problem.

[15] For n=3 and $m_{3/2}$=10TeV, the predicted masses for gauginos in eq.(42) are not excluded by the present experiment [24].

[16] The gravitino loop may also generate gaugino masses as pointed out in Ref. [11]

cut-off scale has turned out to be at the Planck scale. Therefore, it will be interesting to related our discrete R-symmetry to some physics at the Planck scale. A possible candidate is the superstring theory, since it is well known that various discrete (non-R and R-)symmetries arise from compactified manifolds [28] of 10-dimensional space-time. Further investigation on this intriguing possibility will be given elsewhere.

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Figure caption

Fig. 1

Inflaton potentials $V(\phi)$ along the real axis $\mathrm{Re}(\phi)$ are shown for the cases (a) $n = 3$, $v = 1.9 \times 10^{16}$GeV, (b) $n = 4$, $v = 2.5 \times 10^{15}$GeV, (c) $n = 5$, $v = 9.0 \times 10^{14}$GeV. Notice that these sets of parameters are given in Table 1 with $m_{3/2}$=1TeV and the shape of the normalized potentials in this figure are independent of the parameter λ.

$m_{3/2}$=100GeV

n	λ	v[GeV]	Λ[GeV]	m_ϕ[GeV]	T_R[GeV]
3	6.5×10^{-10}	1.1×10^{16}	1.6×10^{25}	2.1×10^{7}	3.8
4	5.0×10^{-7}	1.1×10^{15}	1.6×10^{18}	2.3×10^{9}	4.4×10^{3}
5	1.3×10^{-5}	3.8×10^{14}	1.6×10^{16}	2.5×10^{10}	1.6×10^{5}
6	8.9×10^{-5}	2.0×10^{14}	2.0×10^{15}	1.1×10^{11}	1.4×10^{6}
7	3.1×10^{-4}	1.3×10^{14}	6.6×10^{14}	2.8×10^{11}	5.9×10^{6}
8	7.2×10^{-4}	9.7×10^{13}	3.2×10^{14}	5.6×10^{11}	1.7×10^{7}
9	1.3×10^{-3}	7.9×10^{13}	2.0×10^{14}	9.5×10^{11}	3.7×10^{7}
10	2.1×10^{-3}	6.7×10^{13}	1.5×10^{14}	1.4×10^{12}	6.9×10^{7}

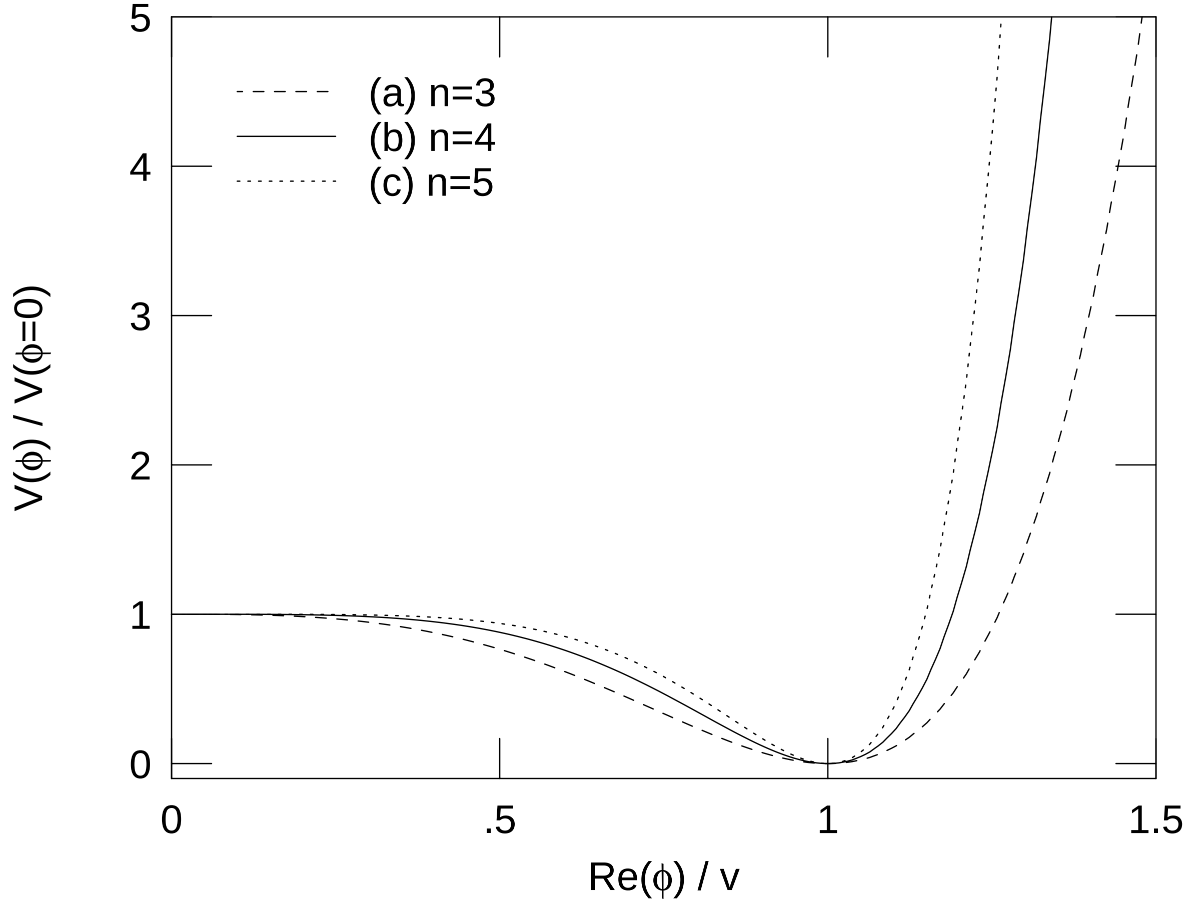
$m_{3/2}$=1TeV

n	λ	v[GeV]	Λ[GeV]	m_ϕ[GeV]	T_R[GeV]
3	1.2×10^{-9}	1.9×10^{16}	1.6×10^{25}	6.6×10^{7}	2.1×10^{1}
4	5.0×10^{-7}	2.5×10^{15}	3.5×10^{18}	4.9×10^{9}	1.4×10^{4}
5	9.9×10^{-6}	9.0×10^{14}	4.2×10^{16}	4.4×10^{10}	3.7×10^{5}
6	5.6×10^{-5}	5.0×10^{14}	5.7×10^{15}	1.7×10^{11}	2.7×10^{6}
7	1.7×10^{-4}	3.4×10^{14}	1.9×10^{15}	4.1×10^{11}	1.1×10^{7}
8	3.7×10^{-4}	2.6×10^{14}	9.7×10^{14}	7.8×10^{11}	2.8×10^{7}
9	6.5×10^{-4}	2.2×10^{14}	6.2×10^{14}	1.3×10^{12}	5.7×10^{7}
10	9.9×10^{-4}	1.9×10^{14}	4.4×10^{14}	1.9×10^{12}	1.0×10^{8}

$m_{3/2}$=10TeV

n	λ	v[GeV]	Λ[GeV]	m_ϕ[GeV]	T_R[GeV]
3	2.1×10^{-9}	3.4×10^{16}	1.6×10^{25}	2.1×10^{8}	1.2×10^{2}
4	5.0×10^{-7}	5.3×10^{15}	7.5×10^{18}	1.1×10^{10}	4.4×10^{4}
5	7.4×10^{-6}	2.1×10^{15}	1.1×10^{17}	7.9×10^{10}	8.8×10^{5}
6	3.5×10^{-5}	1.2×10^{15}	1.6×10^{16}	2.7×10^{11}	5.5×10^{6}
7	9.7×10^{-5}	8.9×10^{14}	5.6×10^{15}	6.0×10^{11}	1.9×10^{7}
8	1.9×10^{-4}	7.0×10^{14}	2.9×10^{15}	1.1×10^{12}	4.5×10^{7}
9	3.2×10^{-4}	5.9×10^{14}	1.9×10^{15}	1.7×10^{12}	8.8×10^{7}
10	4.6×10^{-4}	5.2×10^{14}	1.4×10^{15}	2.4×10^{12}	1.5×10^{8}

Table 1: For gravitino masses $m_{3/2}$=100GeV, 1 and 10TeV, coupling constants λ and vacuum-expectation values v are calculated with $(\delta T/T) = 6.0 \times 10^{-6}$ fixed. We have taken n=3–10, and physical cut-off scales Λ, masses of ϕ field m_ϕ and the reheating temperatures T_R are also shown for each n .



(a) n=3
(b) n=4
(c) n=5
V(ϕ) / V(ϕ=0)
Re(ϕ) / v
0
.5
1
1.5
0
1
2
3
4
5